

Inverse-square law violation and reactor antineutrino anomaly ¹

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Abstract

We discuss a possibility that the so-called reactor antineutrino anomaly can be, at least in part, explained by applying a quantum field-theoretical approach to neutrino oscillations, which in particular predicts a small deviation from the classical inverse-square law at short but macroscopic distances between the neutrino source and detector. An extensive statistical analysis of the reactor data is performed to examine this speculation.

1 Introduction

Nuclear reactors are intense sources of electron antineutrinos whose spectrum is composed of thousands of spectral components formed mainly by the β decay of the fission products of the four parent isotopes: ^{235}U , ^{238}U , ^{239}Pu , and ^{241}Pu . The very sophisticated recent calculations [1–3] yield a net 3–3.5% upward shift in the predicted energy-averaged $\bar{\nu}_e$ flux with respect to the previously expected flux used in the earlier short baseline (SBL) reactor experiments (ILL–Grenoble [4, 5], Gösgen [6], Krasnoyarsk [7, 8], Rovno [9–11], Bugey [12, 13], Savannah River Plant [14]). The $\bar{\nu}_e$ flux normalization uncertainty in the new calculations is claimed to be only $\pm 2.7\%$. This implies [3] that the measured event rates in the SBL experiments are about 6% too low, giving rise to the so-called “reactor antineutrino anomaly” (RAA).

Figure 1 illustrates this issue. The curve shows the ratio of the $\bar{\nu}_e$ -induced event rate calculated with and without regard for the 3ν oscillations. Here and thereafter we use the global best fit values for the neutrino mass-squared splittings and mixing angles from Ref. [15] for the normal mass hierarchy; we also assume no CP violation in mixing. The cross section for the inverse β decay (IBD) is calculated by using the recent analytical results of Ref. [16], which take into account the radiative corrections of order α/π and contributions of weak magnetism and neutron recoil to next-to-leading order in the expansion in inverse powers of the nucleon mass. In Fig. 1 and

¹Talk at the International Workshop on Prospects of Particle Physics: “Neutrino Physics and Astrophysics”, Valday, Russia, February 1–8, 2015.

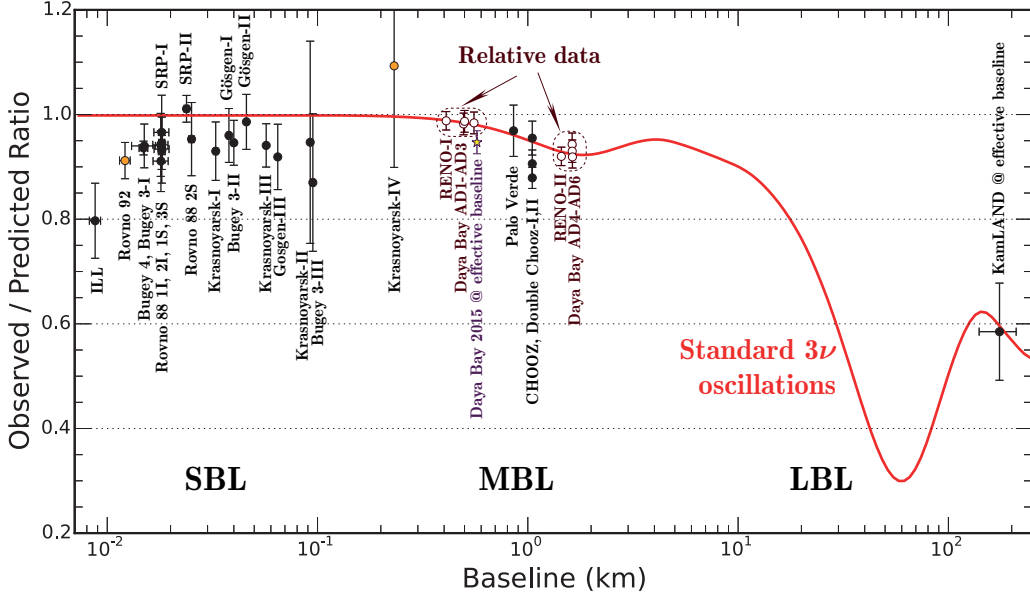


Figure 1: The reactor antineutrino anomaly. The points represent the ratios of measured event rates to that expected with no oscillations. The vertical error bars do not include the common normalization uncertainty. The original data from Refs. [5–14, 17–21] are recalculated as explained in the text. The data from Refs. [22, 23] are relative. The ratios of the KamLAND [21] and Daya Bay 2015 [24] results to the “Huber-Mueller” model predictions are plotted at the flux-weighted distances of 175 km and 573 m, respectively. The curve represents the theoretical expectation, in which the 3ν oscillations parameters are fixed to the global best fit values taken from the Review of Particle Physics [15]; the normal mass ordering and no CP violation are assumed.

below, all the curves correspond to a reactor with pure ^{235}U fuel.² The original data from Refs. [5–14, 17–21] are corrected according to Ref. [1] and then renormalized to the new world average value of the neutron mean life [15]. It is seen that most of the SBL data points (the measurements at $L \lesssim 100$ m, where L is the distance between the reactor core and detector) are below the expectation. A clear trend is visible at $L \lesssim 20$ m that the closer detector is located to reactor, the smaller the measured rate (the larger the discrepancy between the data and theory). Note that the data points “Krasnoyarsk-IV” [8] and “Rovno 92” [11] are ignored in the numerous RAA analyses, but the latter point is significant for revealing the mentioned trend.

²Nevertheless, in the following analysis we explicitly take into account the particular fuel composition in each experiment, since it does have a small impact on the results.

Also shown are the data from the medium and long baseline reactor experiments Palo Verde [17], CHOOZ [18], Double Chooz [19, 20], KamLAND [21], RENO [22], and Daya Bay [23, 24]. The data sets from Refs. [22, 23] are relative measurements, while the recent high-precision Daya Bay measurement [24] is absolute (in Fig. 1 it is placed at the effective baseline of 573 m). As is seen, the theory is in rather poor agreement with the latter result. Hence, both the earlier SBL and new Daya Bay measurements give a hint to either “new physics”, or merely a lower $\bar{\nu}_e$ flux than predicted in Refs. [1, 2].

2 Extra neutrinos or wrong normalization?

Most if not all efforts to resolve the anomaly are based on the hypothesis of existence of one or more light (eV mass scale) sterile neutrinos, that is fundamental neutral fermions with no standard model interactions except those induced by mixing with the standard (active) neutrinos. The active-to-sterile neutrino mixing would lead to a distance-dependent spectral distortion and overall reduction of the reactor $\bar{\nu}_e$ flux.

In Fig. 2 we show, as an example, the results of calculations performed in the framework of the simplest “3+1” phenomenological model with one sterile (anti)neutrino, ν_4 , by using the three pairs of the $\nu_4 - \nu_1$ mixing parameters, $(\Delta m_{41}, \sin^2 2\theta_{41})$, listed in the legend of the figure. These values were derived in Ref. [25] from detailed statistical analyses of all the neutrino oscillation data available to date. The “SBL rates only” fit, includes the SBL reactor data except the points “Krasnoyarsk-IV” and “Rovno 92” (see Fig. 1). The “SBL + Bugey 3 spectrum” fit, includes the same data set and spectral data from Bugey 3 [13]. The “Global ν_e disappearance” fit involves the data from the reactor experiments [4, 6–10, 12–14, 17–19, 22, 23, 26], as well as solar neutrinos (261 data points from Homestake, SAGE, GALLEX/GNO, Super-Kamiokande, and SNO experiments), radioactive source experiments at SAGE and GALLEX, and the LSND/KARMEN ν_e disappearance data from $\nu_e - {}^{12}\text{C}$ scattering (see Ref. [25] for the full list of references and further details). It is necessary to note that these fits operate with somewhat lower (to within roughly 1%) values for the reactor $\bar{\nu}_e$ induced rates and with a little bit different covariance matrix, as compared to those used in the present analysis (the details of our calculations will be published elsewhere).

The solid curve in Fig. 2 represents the same 3ν oscillation prediction as in Fig. 1, but shifted down by the normalization factor N_0 derived from a fit to all the data except these from RENO and Daya Bay. In this fit

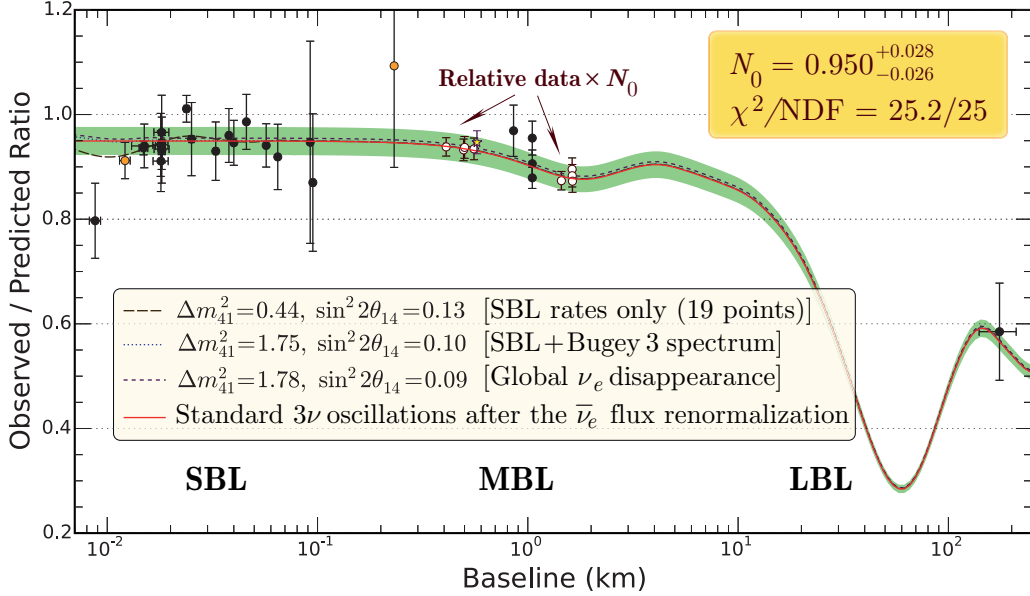


Figure 2: Comparison of the data with the “3+1” model (see text) and with the standard 3ν oscillation prediction after renormalization of the $\bar{\nu}_e$ flux (solid curve). The renormalization factor is obtained from a fit to all the data but RENO and Daya Bay. The filled band shows the $\pm 1\sigma$ uncertainty of the fit. The data points are the same as in Fig. 1 but the eight relative data points are shifted by the factor $N_0 = 0.950$.

we take into account the correlation between the data, including the overall normalization uncertainty, which is taken to be 2.7% [3]. The obtained factor $N_0 = 0.950^{+0.028}_{-0.026}$ ($\chi^2/\text{NDF} \approx 1$) does not contradict to the adopted flux uncertainty, but is somewhat different from the results of earlier calculations [3, 16, 27], which used different data subsets and input parameters. All four curves in Fig. 2 are in agreement, within the errors, with the new Daya Bay measurement [24], but are in some conflict with the ILL data point [5].

Recently, it was argued [28] that the true uncertainty in the $\bar{\nu}_e$ flux predictions may be as large as 5% and the spectral shape uncertainties may be much larger due to poorly known structure of the forbidden decays. This finding has been in essence confirmed by the new precision measurements of the positron energy spectra from IBD [24], which show apparent $\sim 10\%$ excess in 4–6 MeV region compared to the expectation based on the models of Refs. [1, 2] (see Ref. [29] for further discussion and references). From what has been said it appears that the efforts to explain the anomaly by the sterile neutrino hypothesis may be somewhat premature. Moreover, it is seen from

Fig. 2 that the proper renormalization of the flux is hardly distinguishable from the “global ν_e disappearance” 4ν fit and (maybe somewhat accidentally) almost fully coincides with the “SBL + Bugey 3 spectrum” fit. We emphasize however that the steady decrease of the event rate at very small L , *if real*, cannot be explained by neither the wrong $\bar{\nu}_e$ flux normalization alone, nor the “3+1” scenario. Thus, it is appropriate to consider an alternative explanation. Such an alternative has been proposed in Ref. [30]. It is based on a quantum field-theoretical (QFT) approach to the neutrino oscillation phenomenon, which predicts a small deviation of the (anti)neutrino event rate, as a function of the distance L between the source and detector, from the classical inverse-square law (ISL) behavior. Below we consider this issue in some detail.

3 A sketch of the QFT approach

The “neutrino-oscillation” phenomenon in the S -matrix QFT approach is nothing else than a result of interference of the macroscopic Feynman diagrams (like shown in Fig. 3) which describe the lepton number violating

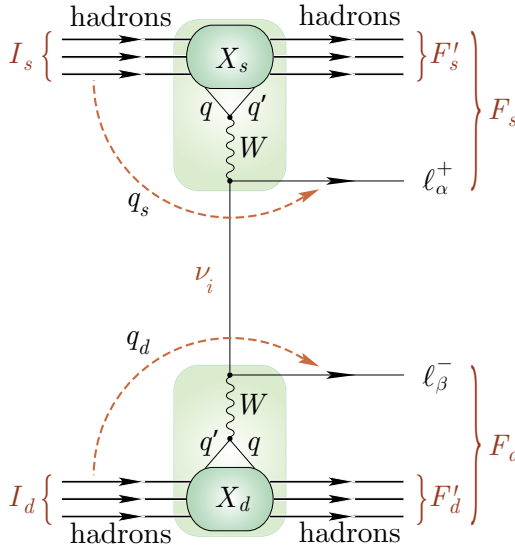


Figure 3: A generic macroscopic Feynman diagram. Here $I_{s,d}$ and $F_{s,d}$ denote the sets of the initial (I) and final (F) WP states in the “source” (X_s) and “detector” (X_d) vertices; $F'_{s,d} = F_{s,d} \oplus \ell_{\alpha,\beta}^{+,-}$, where $\ell_{\alpha,\beta}^{+,-}$ are the charged lepton WP states ($\alpha, \beta = e, \mu, \tau$); $q_{s,d}$ are the 4-momentum transfers in the vertices, as defined by Eq. (5). The vertices are in general *macroscopically* separated in space and time. The particular “decryption” of the neutrino production/absorption mechanism assumes the standard model charged current interaction of quarks and leptons.

processes with the neutrino mass eigenfields ν_i ($i = 1, 2, 3$) treated as internal lines (propagators). The external lines of the macrodiagrams are assumed corresponding to asymptotically free quasi-stable wave packets (WP) rather than the conventional to QFT one-particle Fock’s states $|\mathbf{k}, s\rangle$ with definite

3-momenta \mathbf{k} and spin projections s . According to Refs. [31, 32], the free external WP states are constructed as covariant space-time point x dependent linear superpositions of the one-particle states,

$$|\mathbf{p}, s, x\rangle = \int \frac{d\mathbf{k} \phi(\mathbf{k}, \mathbf{p}) e^{i(k-p)x}}{(2\pi)^3 2E_{\mathbf{k}}} |\mathbf{k}, s\rangle, \quad (1)$$

satisfying the *correspondence principle* which demands that $|\mathbf{p}, s, x\rangle$ turns into $|\mathbf{k}, s\rangle$ in the plane-wave limit (PWL) that is equivalent to the following condition for the relativistic invariant form factor function ϕ :

$$\phi(\mathbf{k}, \mathbf{p}) \xrightarrow{\text{PWL}} (2\pi)^3 2E_{\mathbf{p}} \delta(\mathbf{k} - \mathbf{p}).$$

The detail properties of the WP states (1) are discussed in Refs. [33, 34].

Within the outlined approach, and after applying several more or less “technical” simplifications, it is proved [32] that the neutrino induced event rate in an ideal detector can be written (somewhat symbolically) in the form

$$\frac{dN}{d\tau} = \frac{1}{V_{\mathcal{D}} V_{\mathcal{S}}} \int_{V_{\mathcal{S}}} d\mathbf{x} \int_{V_{\mathcal{D}}} d\mathbf{y} \int d\mathfrak{F}_{\nu} \int d\sigma_{\nu\mathcal{D}} \mathcal{P}_{\alpha\beta}(E_{\nu}, |\mathbf{y} - \mathbf{x}|). \quad (2)$$

Here τ is the detector exposure time, E_{ν} is the neutrino energy, $\mathcal{P}_{\alpha\beta}$ is the QFT generalization of the standard quantum-mechanical neutrino flavor transition probability, the differential form $d\sigma_{\nu\mathcal{D}}$ represents the differential cross section of the neutrino scattering from the whole detector device; $d\mathfrak{F}_{\nu}$ is the differential neutrino flux incident on the detector from a *stationary* source device (e.g., a fission reactor core). The integrations in (2) are over the source and detector fiducial volumes $V_{\mathcal{S}}$ and $V_{\mathcal{D}}$. The theory explicitly predicts that the neutrino flux decreases with increasing the distance $|\mathbf{y} - \mathbf{x}|$ in compliance with the usual inverse-square law (ISL):

$$d\mathfrak{F}_{\nu} \propto |\mathbf{y} - \mathbf{x}|^{-2}. \quad (3)$$

This quite expected result has been derived by using the so-called Grimus-Stockinger (GS) theorem [35], which defines the asymptotic behavior of the amplitude at $L \rightarrow \infty$ and this is the crucial point in the context of the problem under consideration.

As it follows from the formalism, the L dependence of the amplitude described by the macrodiagram shown in Fig. 3 is defined by the neutrino propagator modified by the external wave packets,

$$(i\hat{\partial} + m_i) \int \frac{d^4 q}{(2\pi)^4} \frac{\tilde{\delta}_s(q - q_s) \tilde{\delta}_d(q + q_d) e^{-iqx}}{q^2 - m_i^2 + i\epsilon}, \quad (4)$$

where $x = (y_0 - x_0, \mathbf{y} - \mathbf{x}) \equiv (T, \mathbf{L})$, q_s and q_d are the 4-momentum transfers,

$$q_s = \sum_{a \in I_s} p_a - \sum_{b \in F_s} p_b, \quad q_d = \sum_{a \in I_d} p_a - \sum_{b \in F_d} p_b, \quad (5)$$

p_\varkappa are the most probable (on-shell) 4-momenta of the external packets $\varkappa \in I_s \oplus I_d \oplus F_s \oplus F_d$, and m_i is the mass of the neutrino field ν_i . The functions $\tilde{\delta}_s(q - q_s)$ and $\tilde{\delta}_d(q + q_d)$ are the “smeared” δ functions (see Ref. [30] for their explicit form) defined by the 4-momenta p_\varkappa , masses m_\varkappa ($m_\varkappa^2 = p_\varkappa^2$), and momentum spreads σ_\varkappa of the external in and out packets ($\sigma_\varkappa^2 \lll m_\varkappa^2$). In the plain-wave limit ($\sigma_\varkappa \rightarrow 0, \forall \varkappa$) these functions turn into the ordinary Dirac δ functions, $\tilde{\delta}_s(q - q_s) \xrightarrow{\text{PWL}} \delta(q - q_s)$, $\tilde{\delta}_d(q + q_d) \xrightarrow{\text{PWL}} \delta(q + q_d)$, thus leading to the exact energy-momentum conservation in the vertices of the macrodiagram, and the function (4) becomes, up to a multiplier, the standard fermion propagator. If however the momentum spreads σ_\varkappa are finite, the space-time behavior of the function (4) is nontrivial. In particular, its spatial dependence at sufficiently large distances L is given by the above-mentioned GS theorem [35], according to which³

$$J(\mathbf{L}, \kappa) = \int \frac{d\mathbf{q}}{(2\pi)^3} \frac{\Phi(\mathbf{q}) e^{i\mathbf{q}\mathbf{L}}}{\mathbf{q}^2 - \kappa^2 - i\epsilon} = \frac{e^{i\kappa L} \Phi(-\kappa \mathbf{l})}{4\pi L} \left[1 + \mathcal{O}\left(\frac{1}{\sqrt{L}}\right) \right] \quad \left(\mathbf{l} = \frac{\mathbf{L}}{L} \right)$$

as $L = |\mathbf{L}| \rightarrow \infty$. This offers the QFT explanation of the ISL behavior (3) but does not, however, provide the spatial scale above which the distance L may be considered as “sufficiently large”.

In Ref. [30], an extended version of the GS theorem has been proved, which parametrically defines such a scale by using the asymptotic expansion of the integral $J(\mathbf{L}, \kappa)$ in terms of inverse powers of L at large L . To be more precise, the theorem in its simplest form states that for any function $\Phi(\mathbf{q})$ in the Schwartz space $S(\mathbb{R}^3)$

$$J(\mathbf{L}, \kappa) = \frac{e^{i\kappa L}}{4\pi L} \left[\Phi(\mathbf{q}) + \sum_{n \geq 1} \frac{(-i)^n D_n \Phi(\mathbf{q})}{L^n} \right]_{\mathbf{q} = -\kappa \mathbf{l}}, \quad L \rightarrow \infty, \quad (6)$$

where D_n are explicitly defined differential operators on the momentum space; the lowest order operators, sufficient for our present purpose, are

$$D_1 = \frac{\kappa}{2} [\nabla_{\mathbf{q}}^2 - (\mathbf{l} \nabla_{\mathbf{q}})^2] - (\mathbf{l} \nabla_{\mathbf{q}}),$$

$$D_2 = \frac{\kappa^2}{8} [\nabla_{\mathbf{q}}^2 - (\mathbf{l} \nabla_{\mathbf{q}})^2]^2 - \kappa (\mathbf{l} \nabla_{\mathbf{q}}) [\nabla_{\mathbf{q}}^2 - (\mathbf{l} \nabla_{\mathbf{q}})^2] - \frac{1}{2} [\nabla_{\mathbf{q}}^2 - 3(\mathbf{l} \nabla_{\mathbf{q}})^2].$$

³It is assumed that the complex-valued function $\Phi(\mathbf{q})$ itself and its first and second derivatives decrease at least like $1/|\mathbf{q}|^2$ as $|\mathbf{q}| \rightarrow \infty$ and $\kappa^2 > 0$; see Ref. [30] for details.

Additional important features can be found in Refs. [30, 36]. An analysis of Eq. (6) shows that the $1/L$ behavior of the amplitude (and thus the ISL behavior of the event rate) is violated at the distances $L \lesssim \mathfrak{L}_0$, where

$$\mathfrak{L}_0 \sim \kappa \sigma_{\text{eff}}^{-2} \approx 20 \left(\frac{\kappa}{1 \text{ MeV}} \right) \left(\frac{\sigma_{\text{eff}}}{1 \text{ eV}} \right)^{-2} \text{ cm} \quad (7)$$

and the function $\sigma_{\text{eff}} = \sigma_{\text{eff}}(\kappa; \{\mathbf{v}_\kappa, m_\kappa, \sigma_\kappa\})$ represents an effective momentum spread dependent on the neutrino momentum κ as well as on the mean velocities, masses, and momentum spreads of the external (in and out) wave packets κ . The explicit form of this function can be found after specification of a particular model for the external WP states. A simple example is discussed in Ref. [30] within the so-called contracted relativistic Gaussian packet (CRGP) model [33, 34]. It is in particular shown that σ_{eff} is defined through the transverse (with respect to the neutrino propagation direction \mathbf{l}) components of the inverse overlap tensors which determine the effective space-time overlap volumes of the WP states in the vertices of the macrodiagram. It is significant that these components are nearly independent of the neutrino masses (assuming these to be small with respect to the neutrino energy and thus $\kappa \simeq E_\nu$). Within the CRGP model, it can be also shown that the magnitude of σ_{eff} is strongly affected by the hierarchy of the external momentum spreads σ_κ but in the simplest case when these spreads are similar in order of magnitude, σ_{eff} is of the same order, too. So, as is seen from Eq. (7), the spatial scale (7) can be macroscopically large at sufficiently small external momentum spreads, thus leading to a measurable ISL violation (ISLV).

It is shown in Ref. [33] that Eq. (6) modifies the formula for the event rate (2) in such a way that the relation (3) for the flux is replaced by

$$d\mathfrak{F}_\nu \propto \frac{1}{|\mathbf{y} - \mathbf{x}|^2} \left(1 + \sum_{n \geq 1} \frac{\mathfrak{C}_n}{|\mathbf{y} - \mathbf{x}|^{2n}} \right), \quad (8)$$

where the coefficient functions \mathfrak{C}_n are explicitly defined from Eq. (6). By making expedient assumptions, it can be proved (and this is a crucial point) that $\mathfrak{C}_1 < 0$. Hence, using Eq. (8) in leading order (thereby assuming that the ISLV correction is small), yields the following simple replacement for the event rate:

$$\frac{dN}{d\tau} \longmapsto \left(1 - \frac{\overline{\mathfrak{L}}_0^2}{L^2} \right) \frac{dN}{d\tau} \quad (9)$$

(provided that $\overline{\mathfrak{L}}_0^2 \ll L^2$). Here $\overline{\mathfrak{L}}_0 \sim \langle \mathfrak{L}_0 \rangle$ is a neutrino energy dependent parameter of dimension of length. Needless to say, at present this parameter cannot be obtained from first-principle calculations, but it can be measured.

4 Data analysis

To check the assumption that the ISLV effect could actually be, in part, responsible for RAA, we performed a statistical analysis of the reactor data discussed in Sect. 1. Since in this paper we use only the spectrum-averaged event rates, the L dependent factor in Eq. (9) can be replaced by $(1 - L_0^2/L^2)$, where $L_0 \sim \langle \bar{\mathcal{L}}_0 \rangle$ is an energy independent parameter, which is a subject of the present study. Taking into account the large uncertainty in the $\bar{\nu}_e$ flux normalization, we shall use the following theoretical model to fit the data:

$$T(L; N_0, L_0) = N_0 \left(1 - \frac{L_0^2}{L^2}\right) \frac{\int_0^\infty dE_\nu \sum_k f_k P_{\text{surv}}^{3\nu}(L, E_\nu) \sigma(E_\nu) S_k(E_\nu)}{\int_0^\infty dE_\nu \sum_k f_k \sigma(E_\nu) S_k(E_\nu)}. \quad (10)$$

Here N_0 is the required normalization parameter, f_k is the reactor fissile isotope fraction, $S_k(E_\nu)$ is the $\bar{\nu}_e$ energy spectrum (taken from Ref. [1]), $\sigma(E_\nu)$ is the IBD cross section [16] and $P_{\text{surv}}^{3\nu}(L, E_\nu)$ is the $\bar{\nu}_e$ survival probability in the standard 3ν mixing scheme.⁴ In order to find the best-fit parameters N_0 and L_0 we minimize the standard χ^2 with the full covariance matrix for the correlated data. The results of the fits of several types are shown in Figs. 4–6.

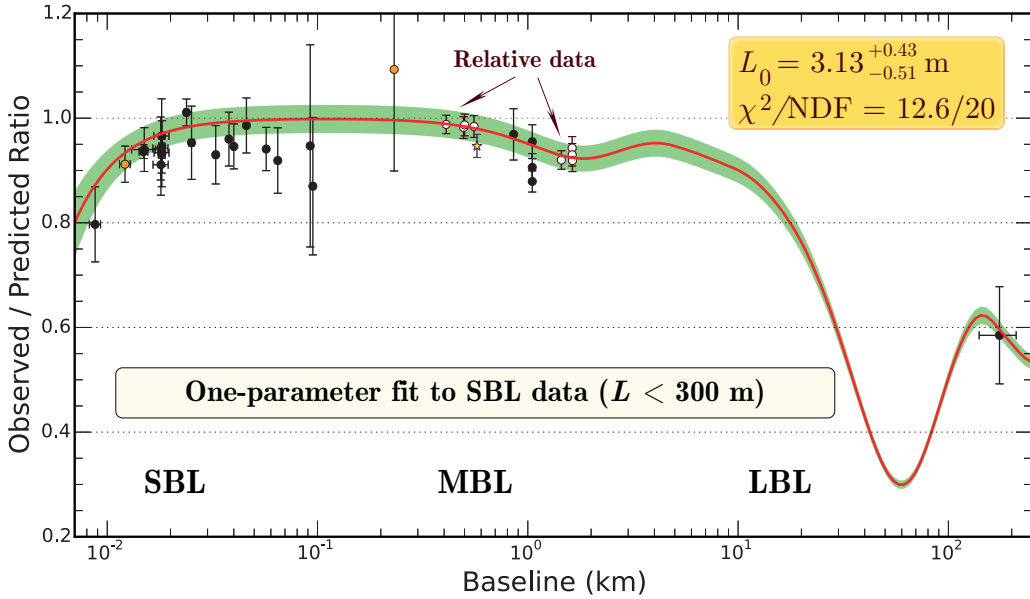


Figure 4: One-parameter fit to SBL data assuming $N_0 = 1$.

⁴We thereby neglect the decoherence effects predicted in the QFT approach [32], reasonably assuming that the baselines under consideration are too short for their manifestation.

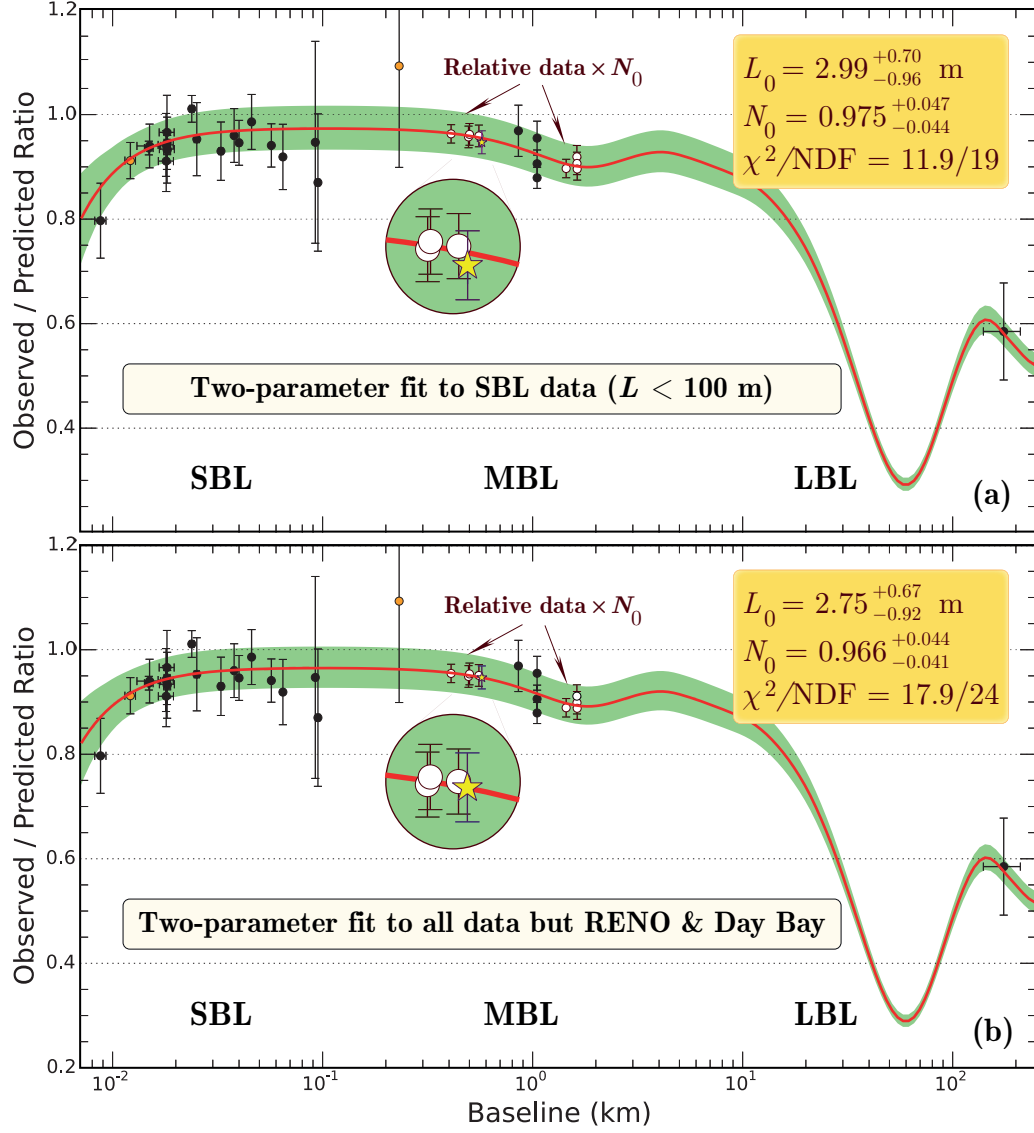


Figure 5: Two-parameter fits to SBL data subset (a) and to all data (b). The data of RENO and Daya Bay are not included into the analysis. The eight relative data points in the panels (a) and (b) are shifted by the factor $N_0 = 0.975$ and 0.966 , respectively. The insets clarify the agreement between the best-fit curves and Daya Bay data points.

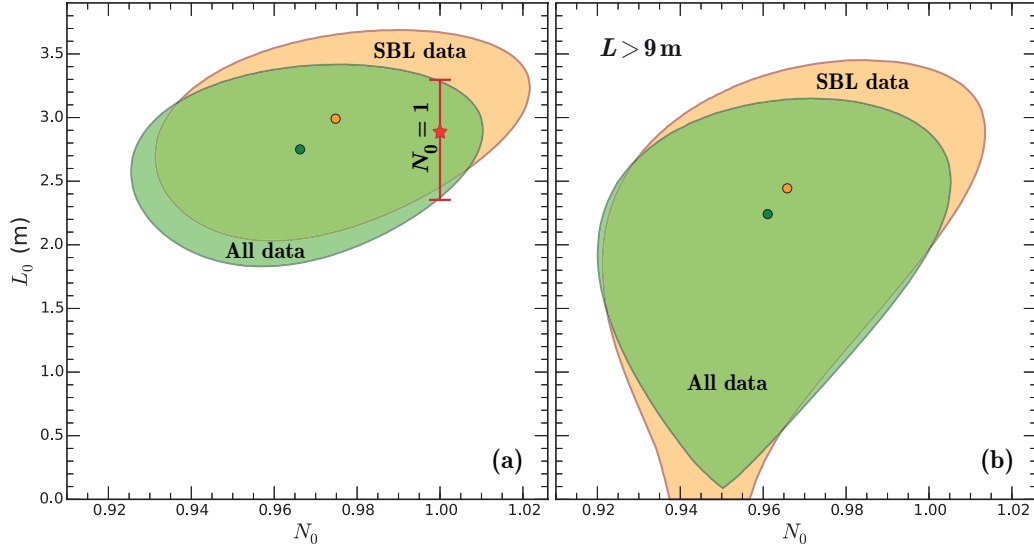


Figure 6: Panel (a): The 68% C.L. error contours for the pair of the parameters N_0 and L_0 extracted from the fits to all data and SBL data subset ($L < 300$ m). The one-parameter (with N_0 fixed to 1) fit to SBL data is also shown for comparison. Panel (b): The same as in (a), but with the ILL data point excluded from the statistical analysis. In both panels, the circles correspond to the best-fit values of the fitted parameters with $L_0^{\text{SBL data}} > L_0^{\text{All data}}$ and $N_0^{\text{SBL data}} > N_0^{\text{All data}}$.

Figure 4 represents the results of the simplest one-parameter fit, in which N_0 is set to 1 and only the SBL data ($L < 300$ m) are used in the analysis. Here and below, the filled band represents the $\pm 1\sigma$ uncertainty. It is seen that despite relatively small value of χ^2/NDF the best-fit curve is in rather poor agreement with the data; in particular, it is in conflict with the recent Daya Bay measurement. This indicates that the ISLV suppression alone is not sufficient and the flux renormalization is actually required.

Figure 5 illustrates the results of the two-parameter fits performed with two different data subsets, namely with the SBL data only, and with all the data except these from RENO and Daya Bay. As is seen from the figure, these two fits are in reasonable agreement to each other and both describe the reactor data rather well. Although the value of χ^2/NDF for the SBL data is nominally a bit better than that for the fit to all data (0.63 and 0.74, respectively), the latter fit is (as is clearly seen in the inserts of Fig. 5) in better agreement with the Daya Bay point (let us remind that it does not participate in the analysis). By comparing these results with the fits shown in Fig. 2, we may conclude that the ISLV effect in combination with the proper renormalization of the $\bar{\nu}_e$ flux provides a better resolution of the anomaly.

To gain a deeper understanding of our results, we compare in Fig. 6 the 68% C.L. error contours for the pair of the fitted parameters (N_0, L_0) , obtained from the fits to different data subsets. Panel (a) in Fig. 6 shows the contours for the fits to all data and SBL data subset. Panel (b) shows the same but with the ILL data point excluded from the analysis. It is seen with no need of additional explanation that the Grenoble experiment is the “cornerstone” for verification of the ISLV effect. It is however important that even without the ILL point we obtain essentially the same best-fit values of the parameters.

5 Conclusions

The QFT approach predicts a deviation from the classical inverse-square law at short baselines. While the numerical value of the spatial scale at which the deviation becomes essential cannot be predicted from the present-day theory, it can be extracted, under reasonable assumptions, from the data of the past and current reactor antineutrino experiments. Our statistical analysis demonstrates that the averaged over the reactor antineutrino spectrum value of the scale (L_0) is about 3 m that roughly corresponds to the spectrum-averaged effective momentum spread $\langle\sigma_{\text{eff}}\rangle$ of about $0.5 - 0.8$ eV (thereby hinting that the wave packets of the particles and nuclei involved into the reactor $\bar{\nu}_e$ production and detection may have “mesoscopic” effective dimensions). This is in agreement with the conservative estimate presented in Ref. [30]. Besides, the best-fit value of L_0 is very stable with respect to choice of the data subset and $\bar{\nu}_e$ spectrum model. To check the latter, we performed the same one- and two-parameter fits as described above, but with the input $\bar{\nu}_e$ energy spectra from Refs. [37] and [38] derived by very different methods, as well as with combinations of the models [1, 2] and the cumulative $\bar{\nu}_e$ spectrum from ^{238}U fission recently measured with the scientific neutron source FRM II in Garching [39]. We conclude from these exercises that value of L_0 is almost insensitive (within the errors) to the spectrum variations. Needless to say, it is not the case for the normalization parameter N_0 .

Although the available reactor data cannot definitely confirm or exclude the light sterile neutrino hypothesis, and do not provide unambiguous support for the ISLV effect, they are in much better agreement with the latter. The next-generation experiments with very short baselines ($L \lesssim 20$ m), small neutrino or antineutrino sources, and high-precision, desirably movable detectors are required in order to confirm or disconfirm our explanation.

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